

# A Shift Symmetry in the Higgs Sector: Experimental Hints and Stringy Realizations

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## Abstract

We interpret reported hints of a Standard Model Higgs boson at  $\sim 125$  GeV in terms of high-scale supersymmetry breaking with a shift symmetry in the Higgs sector. More specifically, the Higgs mass range suggested by recent LHC data extrapolates, within the (non-supersymmetric) Standard Model, to a vanishing quartic Higgs coupling at a UV scale between  $10^6$  and  $10^{18}$  GeV. Such a small value of  $\lambda$  can be understood in terms of models with high-scale SUSY breaking if the Kähler potential possesses a shift symmetry, i.e., if it depends on  $H_u$  and  $H_d$  only in the combination  $(H_u + \overline{H}_d)$ . This symmetry is known to arise rather naturally in certain heterotic compactifications. We suggest that such a structure of the Higgs Kähler potential is common in a wider class of string constructions, including intersecting D7- and D6-brane models and their extensions to F-theory or M-theory. The latest LHC data may thus be interpreted as hinting to a particular class of compactifications which possess this shift symmetry.

# 1 Introduction and Summary

While the LHC has so far not produced any significant hint of low-scale supersymmetry, signals of a Standard Model (SM) Higgs boson at  $\sim 125$  GeV have been reported [1]. It is clearly far too early to give up on TeV-scale SUSY (see e.g. [2]). Nevertheless, it may be worthwhile investigating what the latest data, if substantiated, imply for string model building without low-energy supersymmetry. Even if low-energy supersymmetry were not to be found at the LHC, it is out of question that string theory remains the most successful candidate for an ultra-violet completion of gauge and gravitational interactions. Let us therefore search for possible interpretations of the announcements [1] within string theory without assuming an imminent discovery of supersymmetry.<sup>1</sup>

It has been known for a long time [4–7] that, within the SM, certain values of the Higgs mass relate to a vanishing quartic coupling  $\lambda$  at some high but sub-Planckian energy scale. This includes some predictions of a  $\sim 125$  GeV Higgs [5, 7] and has been further discussed in a number of recent papers [8, 9]. We illustrate the situation in figure 1. In the following, we will be concerned with possible origins of the apparently favorable high-scale boundary condition  $\lambda = 0$ .

The necessary fine-tuning of the electroweak scale within the non-supersymmetric SM can arise for example within the flux-based string theory landscape (see e.g. [10]). The best-understood underlying string compactifications are nonetheless supersymmetric at least near the string scale. Hence we assume that at some high energy scale the SM is embedded in a supersymmetric theory. This puts us in the context of *high-scale* SUSY breaking (see e.g. [11–13]), where we can start from the tree-level expression

$$\lambda(m_S) = \frac{g^2(m_S) + g'^2(m_S)}{8} \cos^2 2\beta \quad (1)$$

for the SM quartic coupling at the SUSY breaking scale  $m_S$  with  $g$  and  $g'$  the gauge couplings of  $SU(2)_L$  and  $U(1)_Y$ , respectively. The desire to have vanishing  $\lambda$  at some high energy scale thus points towards models with  $\tan \beta \simeq 1$ . The interesting question is now which structure of the high-scale model may be responsible for this particular value of  $\beta$ .

A possible answer can be based on a rather old observation in the heterotic orbifold context [16–19] (also investigated more recently in the orbifold GUT context [20–23]). The observation is that, for a certain class of models, the Higgs-sector Kähler potential possesses a shift symmetry,  $H_u \rightarrow H_u + c$ ,  $H_d \rightarrow H_d - \bar{c}$ , at tree-level. It hence reads

$$K_H = K_H(H_u + \bar{H}_d, \bar{H}_u + H_d, S, \bar{S}) = |H_u + \bar{H}_d|^2 f(S, \bar{S}) + \dots, \quad (2)$$

where  $S$  stands for an appropriate set of moduli. In section 2 we will show that this structure indeed corresponds to  $\lambda = 0$  at the supersymmetry breaking scale. We will furthermore analyse the correlation between the value of the Higgs mass suggested by [1] and the supersymmetry breaking scale  $m_S$  under the assumption of this shift symmetry.

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<sup>1</sup> For an alternative approach see e.g. the recent analysis of [3].

Interpreted as a prediction for  $m_S$  in terms of the Higgs mass, the concrete result strongly depends on the precise value of the top-quark mass and  $\alpha_s$ . Note that this uncertainty affects not only our approach, but is in fact common to all predictions of the Higgs mass on the basis of RG running arguments, including e.g. the analyses of [5, 7]. We also quantify loop corrections to the shift-symmetric tree-level Kähler potential from the Yukawa couplings. We argue that these loop effects can be treated as a small perturbation and estimate their impact on Higgs mass predictions.

Our philosophy in this letter is to take the phenomenological considerations of section 2 as a motivation to investigate, in section 3, how a shift symmetry of type (2) follows from ultra-violet completions of the SM within string theory. We will argue that a shift symmetry of the above type can be realised in string models where the Higgs sector is related to Wilson line moduli of a higher dimensional gauge theory<sup>2</sup>. As noted already, examples of this type of bulk Higgs fields and the associated shift symmetries are known to arise in orbifold compactifications of the heterotic string [16–19]. We will argue that similar structures are possible in suitable Type II compactifications with D-branes, thereby taking some first steps towards generalising the framework of shift-symmetric Higgs sectors beyond the heterotic orbifold context.

As a starting point it is beneficial to recall the geometric or field theoretic origin of the shift symmetry within the orbifold constructions of [16, 17, 19] as discussed in [22]. We are interested in situations where both Higgs doublets come from the untwisted sector, more specifically from gauge field components associated with the same complex plane. Such fields are known as continuous Wilson lines [24]. They may be defined in a manifestly gauge-invariant manner by the holonomy corresponding to a certain closed loop in the orbifold (obviously not corresponding to an element of  $H_1$ , but impossible to contract because of the orbifold singularities). One may also think of the physical meaning of these continuous Wilson lines in terms of the relative orientation in gauge-space of subgroups surviving at distant fixed-points (see section 4 of [25]).

Translated into the language of smooth heterotic models [28], such fields will contribute to certain types of bundle moduli.<sup>3</sup> These in turn should correspond to brane moduli in Type II orientifolds with D-branes. In particular, it can be the Wilson line sector counted by  $H^1(\Sigma)$  of a D-brane wrapping a cycle  $\Sigma$  of the compactification space in which Higgs fields with a shift symmetry take their origin.

The presence of a shift symmetry in the tree-level Kähler potential and its absence in the superpotential can be understood in field theoretic terms as follows: A non-zero vacuum expectation value (VEV) of a Wilson line cannot be detected by any local observer in the original higher-dimensional theory. Thus, tree-level dimensional reduction at quadratic order does not see this VEV. However, at cubic order (corresponding e.g. to Yukawa couplings coming from gauge couplings) the zero modes of matter fields couple to the Wilson line and hence to the Higgs. This is clear since a zero-mode is a global object and it potentially feels the holonomy along some closed loop on the compact space. Thus,

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<sup>2</sup>In this sense our proposal can be viewed a stringy/supersymmetric version of [5].

<sup>3</sup>In addition, the bundle moduli also receive contributions from twisted sector fields in the orbifold limit, see e.g. [29].

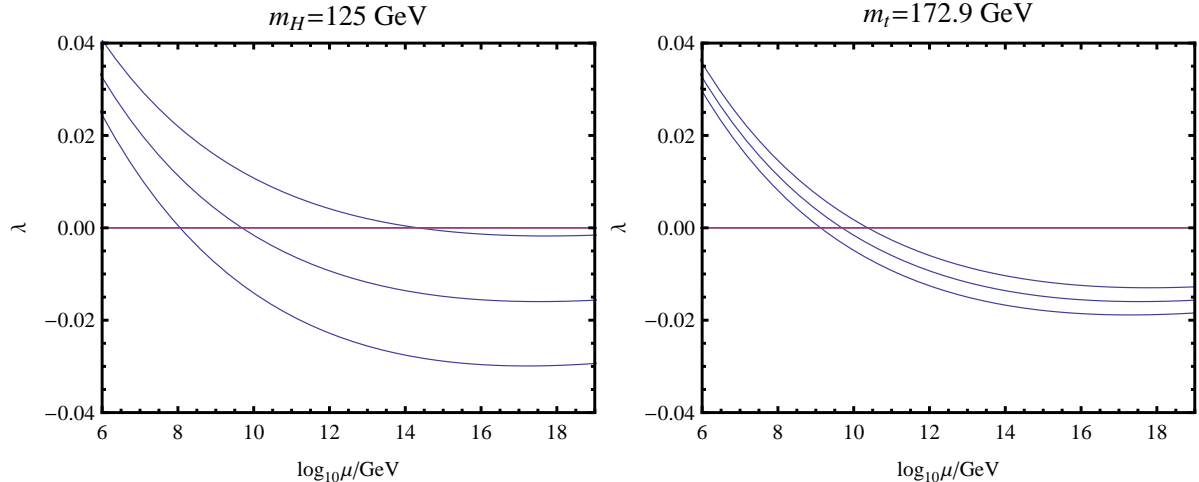


Figure 1: The two-loop RG running of the Higgs quartic coupling in the SM. Left:  $m_H = 125$  GeV,  $m_t = 170.7, 172.9, 175$  GeV from top to bottom. Right:  $m_t = 172.9$  GeV,  $m_H = 126, 125, 124$  GeV from top to bottom.

both in smooth heterotic models and in other string compactifications (e.g. with gauge theories from branes) we expect analogous shift symmetries to arise whenever we have a bundle/brane deformation which is a *pure gauge transformation* in the (sufficiently small) neighbourhood of any point in the higher-dimensional compact space. Its physical reality as a true deformation must be associated purely with non-local effects, as is the case for the famous heterotic continuous Wilson line. Furthermore, since we need only an approximate shift symmetry phenomenologically, we may be satisfied with models which fulfill the above requirement only approximately. In section 3 we will back up these general considerations applied to Type II compactifications with D-branes by a well-known conformal field theory argument [33, 34].

The non-trivial task is clearly to construct models where (part of) the MSSM Higgs degrees of freedom are realized in this way, with a sizeable top Yukawa coupling still present. It is known from heterotic orbifold constructions with Gauge-Yukawa-unification [27] that this is possible in principle. In section 3 we make some preliminary steps towards generalising this structure to Type II compactifications with D-branes. While the appearance of a Higgs sector from Wilson lines in the above sense seems natural, a detailed investigation of realistic models is an exciting challenge left for future work.

## 2 Phenomenology of Higgs sector shift symmetry

We begin our phenomenological analysis by reviewing in more detail how the high-scale boundary conditions for the SM quartic coupling arise in the four-dimensional supergravity picture. We then demonstrate that a shift symmetry of the Kähler potential

at high scales is a predictive assumption even in the presence of the top Yukawa coupling which violates it at the one-loop level.

Our starting point is the Kähler potential (2), and we assume that  $W$  contains no term  $\sim H_u H_d$  (i.e. the  $\mu$  term is generated solely through the Giudice-Masiero mechanism). Without loss of generality, we take  $f = 1$  in the vacuum. Upon supersymmetry breaking, soft masses  $m_{H_u}^2, m_{H_d}^2$ , the  $B\mu$  term and the effective  $\mu$  term are generated. The resulting Higgs mass matrix (with  $B\mu \equiv m_3^2$ ), defined through

$$\mathcal{L} \supset -m_1^2 |H_u|^2 - m_2^2 |H_d|^2 - m_3^2 (H_u \bar{H}_d + \bar{H}_u H_d), \quad (3)$$

is then given by (cf. e.g. [19])

$$m_1^2 = m_2^2 = m_3^2 = |\mu|^2 + m_{3/2}^2 - F^S \bar{F}^{\bar{S}} (\ln f)_{S\bar{S}}, \quad (4)$$

where

$$|\mu|^2 = \left| m_{3/2} - \bar{F}^{\bar{S}} f_{\bar{S}} \right|^2, \quad \bar{F}^{\bar{S}} = e^{K/2} K^{\bar{S}S} D_S W \quad \text{and} \quad m_{3/2} = e^{K/2} W. \quad (5)$$

The generalization to several moduli instead of  $S$  is obvious. The Higgs mass matrix (4), which owes its special form  $m_1^2 = m_2^2 = m_3^2$  to the shift-symmetric Kähler potential (2), has the peculiar property that there is one vanishing eigenvalue with the eigenvector

$$H_0 = \frac{1}{\sqrt{2}} (H_u - \bar{H}_d). \quad (6)$$

Since we assume that the soft scale is at least several orders of magnitude above the electroweak scale, we are thus in the decoupling limit where one SM-like Higgs doublet, the field  $H_0$ , remains light and provides the SM Higgs boson and the would-be Goldstone modes for  $W^\pm$  and  $Z$ . The orthogonal combination and thus the states corresponding to the heavy and charged Higgs as well as the pseudoscalar Higgs in the MSSM become heavy at the soft scale. The SM Higgs boson  $H_0$  as in (6) corresponds to a Higgs mixing angle  $\tan \alpha = -1$ . Since the electroweak symmetry breaking VEV resides only in the light Higgs boson in the decoupling limit, this also implies  $\tan \beta = \tan(\alpha + \frac{\pi}{2}) = 1$ . As a consequence, the tree-level quartic coupling for  $|H_0|^4$  originating from the D-term potential of the electroweak gauge theory vanishes at the soft scale  $m_S$  according to (1). Furthermore, in this scenario the gauge and Yukawa couplings to the light Higgs boson have exactly their respective SM values. This means for example that, unlike in MSSM-like scenarios with large  $\tan \beta$ , the top quark has the only  $\mathcal{O}(1)$  Yukawa coupling.

Let us say a few more words concerning the phenomenological implications. We can base ourselves completely on the analyses of e.g. [9, 12], which give an explicit formula for the SM Higgs mass based on linearized approximations of the RGE solutions. Adapted to our case of interest, they imply (with all mass scales in GeV)

$$m_h = 125 + 1.0 \left( \log_{10} \frac{m_S}{4 \cdot 10^9} \right) + 1.8 \left( \frac{m_t - 173.2}{0.9} \right) - 0.5 \left( \frac{\alpha_s(m_Z) - 0.1184}{0.0007} \right) + \delta. \quad (7)$$

This simply states that  $\tan \beta = 1$  (i.e.  $\lambda = 0$ ) at the SUSY breaking scale  $m_S = 4 \cdot 10^9$  GeV implies  $m_h = 125$  GeV with a certain set of corrections given in self-explanatory notation. Obviously, this has the potential of leading to a rather precise prediction of the SUSY breaking scale. We emphasize, however, that such a prediction requires a significant improvement of the top and Higgs mass measurements: At present, a  $2\sigma$  shift of  $m_t$  together with a 126 GeV Higgs mass allows one to move the point where  $\lambda$  vanishes all the way up to the Planck scale [9].<sup>4</sup>

The crucial issue for us is the intrinsic theoretical uncertainty  $\delta$  in Eq. (7) associated with violations of the shift symmetry. Let us estimate how small this uncertainty might become under favorable circumstances: We take the Kähler potential to be exactly shift-symmetric at the compactification scale  $m_C$ . However, the shift symmetry is broken by the Yukawa couplings in the superpotential. This will feed into the Kähler potential at the very least through field-theoretic loops in the energy range  $m_C > E > m_S$ . We may hope to be in a setting where  $m_S$  and  $m_C$  are closely related, but an exact equality is hard to imagine (or even to properly define). We thus need to estimate the magnitude of these shift-symmetry-violating corrections to the Kähler potential in the supersymmetric theory.

We expect the leading contributions to the Kähler potential to be those of the rigid SUSY limit. These are proportional to  $H_u \bar{H}_u$  and  $\bar{H}_d H_d$  from Higgs self-energy graphs involving gauge and Yukawa couplings. We can thus write the one-loop contribution as ( $t = \ln Q/Q_0$ )

$$\frac{d}{dt}K \sim -2\gamma_{H_u}(S, \bar{S}) H_u \bar{H}_u - 2\gamma_{H_d}(S, \bar{S}) \bar{H}_d H_d. \quad (8)$$

The corresponding one-loop anomalous dimensions in the MSSM are

$$\gamma_{H_u} \Big|_{\theta=\bar{\theta}=0} \sim \frac{3}{16\pi^2} \left( |y_t|^2 - \frac{1}{2}g_2^2 - \frac{1}{10}g_1^2 \right), \quad \gamma_{H_d} \Big|_{\theta=\bar{\theta}=0} \sim \frac{3}{16\pi^2} \left( -\frac{1}{2}g_2^2 - \frac{1}{10}g_1^2 \right). \quad (9)$$

The coupling constants are given by  $y_t = \sqrt{2}y_t^{SM}$  for  $\tan \beta = 1$ ,  $g_2 = g$  and  $g_1 = \sqrt{5/3}g'$  in terms of the usual SM quantities. In the presence of supersymmetry breaking, the resulting running of the soft parameters is captured by the moduli dependence of the anomalous dimensions [14, 15]. The latter can be inferred from the moduli dependence of gauge and Yukawa couplings and Z-factors according to

$$\gamma_{H_u}(S, \bar{S}) = \frac{3}{16\pi^2} \left( \frac{|y_t(S)|^2}{Z_{Q_3}(S, \bar{S})Z_{u_3}(S, \bar{S})} - Z_{H_u}(S, \bar{S}) \left( \frac{g_2^2(S, \bar{S})}{2} + \frac{g_1^2(S, \bar{S})}{10} \right) \right) \quad (10)$$

and analogously for  $\gamma_{H_d}$ . In the rigid SUSY limit, the corrections to the Kähler potential (8) give rise to the well-known running of soft terms and  $\mu$ . Thus, we can equivalently

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<sup>4</sup>This is not well reflected by the *linearized* expression in Eq. (7). We numerically solve the full two-loop RGEs in the following.

work with MSSM one-loop RGEs [30]

$$\begin{aligned}
16\pi^2 \frac{d\mu}{dt} &\sim 3\mu \left( |y_t|^2 - g_2^2 - \frac{1}{5}g_1^2 \right), \\
16\pi^2 \frac{dB\mu}{dt} &\sim 3B\mu \left( |y_t|^2 - g_2^2 - \frac{1}{5}g_1^2 \right) + 6\mu \left( |y_t|^2 A_t + g_2^2 M_2 + \frac{1}{5}g_1^2 M_1 \right), \\
16\pi^2 \frac{dm_{H_u}^2}{dt} &\sim 6|y_t|^2 (m_{H_u}^2 + m_{Q_3}^2 + m_{u_3^c}^2) + 6|y_t|^2 |A_t|^2 - 6g_2^2 |M_2|^2 - \frac{6}{5}g_1^2 |M_1|^2, \\
16\pi^2 \frac{dm_{H_d}^2}{dt} &\sim -6g_2^2 |M_2|^2 - \frac{6}{5}g_1^2 |M_1|^2.
\end{aligned} \tag{11}$$

For  $M_1 = M_2$ , the dependence of these RGEs on the dimensionless parameters can be expressed entirely through the functions  $\beta_y = 6|y_t|^2/16\pi^2$  and  $\beta_g = 3(-g_2^2 - \frac{1}{5}g_1^2)/16\pi^2$ . It is therefore useful to define

$$\epsilon_{y,g} \equiv \int_{\ln m_C}^{\ln m_S} \beta_{y,g}(t) dt \tag{12}$$

as the small parameters controlling the corrections. We now want to estimate the impact of shift-symmetry-violating interactions on  $\tan \beta$  in this setting. As discussed above, we evaluate the soft parameters according to (4) at the scale  $m_C$  and use their conventional renormalization group evolution down to the scale  $m_S$  giving us the mass matrix  $m_1^2(m_S) \dots m_3^2(m_S)$ . The resulting mass matrix is generically nonsingular, and we are (un-surprisingly) confronted with the Higgs hierarchy problem: a finely tuned contribution to the Higgs potential (which we expect to be similar in size to the radiative corrections above) must arise in order to have  $v \ll m_S$ . The resulting massless state is then given by

$$H_0 \sim \frac{1}{\sqrt{2}} \left( -1 \pm \frac{1}{2} \cos 2\beta \right) H_u + \frac{1}{\sqrt{2}} \left( 1 \pm \frac{1}{2} \cos 2\beta \right) \bar{H}_d \tag{13}$$

where the value of  $\cos 2\beta \ll 1$  depends on the exact form of the correction which tunes the electroweak scale small. We can now consider two types of corrections: for small hierarchies with  $|\ln m_S/m_C| \lesssim 1$ , further contributions to the mass matrix beyond those considered here can easily arise and tune the electroweak scale small. For larger hierarchies  $m_S \ll m_C$ , it appears more consistent to tune the weak scale entirely within a leading-log approximation: We tune mass parameters  $\mu$ ,  $M_i$ ,  $m_Q^2$ ,  $m_u^2$  and the trilinear coupling  $A_t$  such that the leading-log corrected matrix defined by (11),

$$M(m_S) = (|\mu|^2 + m_H^2) \begin{bmatrix} 1 & 1 \\ 1 & 1 \end{bmatrix} + \begin{bmatrix} \delta|\mu|^2 + \delta m_{H_u}^2 & \delta B\mu \\ \delta B\mu & \delta|\mu|^2 + \delta m_{H_d}^2 \end{bmatrix}, \tag{14}$$

becomes singular. To leading order in  $\epsilon$ , the condition for vanishing determinant of the scalar mass matrix at the soft scale is

$$2\delta|\mu|^2 + \delta m_{H_u}^2 + \delta m_{H_d}^2 - 2\delta B\mu = 0, \tag{15}$$

which for universal stop, higgs and electroweak gaugino soft masses  $m_{\tilde{q}}^2$ ,  $m_H^2$  and  $M_{1/2}$  corresponds to

$$[2m_{\tilde{t}}^2 + (A_t - \mu)^2] \epsilon_y + 2 [2M_{1/2}^2 - m_H^2 + 2\mu M_{1/2} + \mu^2] \epsilon_g = 0. \tag{16}$$

Assuming that the condition (15) is satisfied, the corrected mass matrix to first order in  $\epsilon_i$  has by assumption one eigenvector with eigenvalue  $\mathcal{O}(\epsilon^2)$  which approximately corresponds to the light Higgs state. It is given by

$$\left(-1 + \frac{1}{2} \frac{\delta m_{H_u}^2 - \delta m_{H_d}^2}{|\mu|^2 + m_H^2}, 1\right) + \mathcal{O}(\epsilon^2) \quad (17)$$

and by comparing with (13) we can directly read off

$$\cos 2\beta = \epsilon_y \frac{m_H^2 + m_{Q_3}^2 + m_{u_3}^2 + |A_t|^2}{2(|\mu|^2 + m_H^2)} \quad (18)$$

and thus  $\cos^2 2\beta \sim \epsilon_y^2$  up to  $\mathcal{O}(1)$  factors. The effective quartic coupling at the soft breaking scale, up to other SUSY thresholds [12], is then given by

$$\lambda(m_S) = \frac{1}{8} (g^2(m_S) + g'^2(m_S)) \cos^2 2\beta. \quad (19)$$

The impact of these shift-symmetry-violating effects on the Higgs mass is illustrated in Figure 2 using the SM two-loop RGEs for the quartic, Yukawa and gauge couplings. For the sake of concreteness, we plot the cases  $m_C^2 = m_S M_{Pl}$  and  $m_S = 10^{-2} m_C$  for  $\cos^2 2\beta = 0 \dots 2\epsilon_y^2$ .

We can conclude that the tree level shift symmetry in the Kähler potential has predictive power despite the presence of the top Yukawa coupling as long as some reasonable assumptions about the nature of SUSY breaking are made. Due to the large impact of the top Yukawa coupling on the running of the quartic coupling, our prediction of the soft scale  $m_S$  for  $m_h = 124 \dots 125$  GeV varies between  $m_S \sim 10^6 \dots 10^{19}$  GeV. This situation will of course improve with more precise knowledge of the top quark and Higgs boson mass. The soft scales and compactification scales which appear in our model for intermediate values of the top quark mass are in a favourable range for neutrino mass generation via the seesaw mechanism and, if they are related to  $f_a$ , to make the axion a realistic dark matter candidate. Note that for one of the parameter points satisfying  $m_C^2 = m_S M_{Pl}$ , namely  $m_S \approx 10^9$  GeV and  $m_C \approx 10^{14}$  GeV, the coupling constants of  $SU(2)_L$  and  $U(1)_Y$  unify at the compactification scale  $m_C$  in standard GUT normalization. This might be interesting in the context of models with symmetry breaking patterns of the type discussed in this work. Of course, these relations can vary depending on the precise embedding of the Higgs field into the adjoint representation of the relevant gauge group.

### 3 Towards a stringy model

In this section we argue that in certain string models a class of bulk matter enjoys a leading-order shift symmetry in the low-energy effective action that gives rise to a Higgs Kähler potential of the type (2). For this to happen the bulk matter must be related to Wilson line moduli, whose typical shift symmetry is then responsible for the advocated



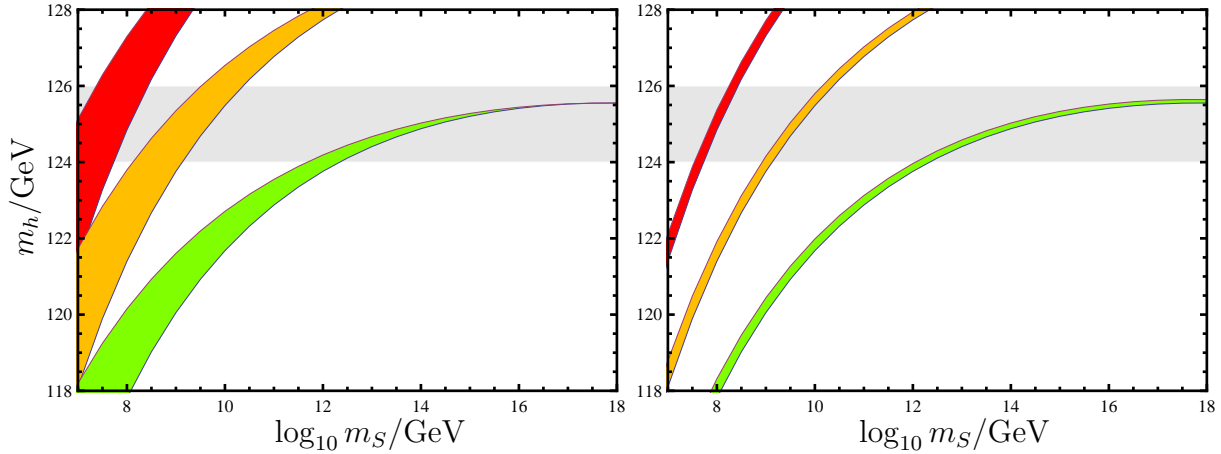


Figure 2: The Higgs mass as a function of the top mass and the soft breaking scale for  $\cos^2 2\beta \in [0 \dots 2\epsilon_y^2]$ . Shown here are the case of a variable compactification scale  $m_S = m_C^2/M_{Pl}$  (left) and  $m_S = 10^{-2}m_C$  (right). The lower (green), middle (orange) and upper (red) bands correspond to top masses of  $m_t = 170.7, 172.9$  and  $175$  GeV respectively. The strong coupling is fixed at  $\alpha_s(m_Z) = 0.1184$ . A shaded band  $m_h = 124 \dots 126$  GeV is included for orientation.

structure. While originally observed in the context of heterotic orbifold models, this phenomenon is much more general and includes Type II orientifold models with bulk matter along D-branes.

### 3.1 Shift symmetry in heterotic orbifold models

A Kähler potential of the structure (2) for the Higgs field arises in heterotic orbifold models where the Higgs fields emerge from states in the untwisted sector [16–18]. Such excitations are the remnants of the internal polarisation states associated with the original ten-dimensional  $E_8 \times E_8$  (or  $\text{Spin}(32)/\mathbb{Z}_2$ ) vector multiplet after imposing the orbifold projection. They therefore propagate as “bulk matter” on the internal torus orbifold. By contrast, states from the so-called twisted sector, which are not inherited from the original vector multiplet, do in general not exhibit a Kähler potential of the form (2). Such states are present only after taking the orbifold quotient, and satisfy twisted boundary conditions. They are localised at the orbifold fixed-points.

Specifically, [16] has analysed  $\mathcal{N} = 1$  supersymmetric  $(0, 2)$   $\mathbb{Z}_N$  heterotic orbifolds in which the internal six-torus factorises as  $T^6 = T^4 \times T^2$ . The Kähler potential associated with moduli of the  $T^2$  factor takes the form

$$K = -\ln[(T + \bar{T})(U + \bar{U}) - (B + \bar{C})(\bar{B} + C)], \quad (20)$$

where  $T$  and  $U$  represent the Kähler and complex structure modulus of the  $T^2$ . This was derived in the supergravity analysis of [16] as a consequence of the coset structure of the

moduli space characteristic for such models. The role of the Higgs fields  $H_u$  and  $H_d$  in (2) is played by the complex Wilson line moduli  $B$  and  $C$  arising as combinations of the the Wilson lines of the gauge field along the two one-cycles of  $T^2$ . These supergravity results agree with the CFT analysis of [17, 26], where the Kähler potential of heterotic (2, 2) models has been computed to second order in momenta by explicit evaluation of 4-point scattering amplitudes. Indeed the above structure of the Kähler potential was confirmed, where the role of  $B$  and  $C$  is played by chiral superfields in representation **27** and  $\overline{\mathbf{27}}$  of  $E_6$  as long as these emerge from the untwisted sector. The shift symmetry receives corrections at one-loop order in the string coupling  $g_s$ , as computed in [17]. Furthermore, there are possibly corrections from couplings to twisted sector fields encoded in higher order (but at treelevel in  $g_s$ )  $n$ -point amplitudes. As pointed out in [17], if the gauge group is further broken e.g. by discrete Wilson lines as in the Hosotani-Witten mechanism, the descendents of these fields continue to exhibit the desired structure in the leading-order tree-level Kähler potential. Implications of the Kähler potential (20) for generation of a  $\mu$ -term and related phenomenological aspects have been analysed in [18].

Interestingly a Higgs sector emerging from untwisted matter in the bulk of the heterotic orbifold is a characteristic of more recent, realistic heterotic orbifold model building. This structure offers, among other things, a solution to the doublet-triplet splitting problem. For constructions of this type and further references see e.g. [27].

The origin of untwisted matter states as descendents of the ten-dimensional gauge multiplet suggests an interpretation of the peculiar structure of the Kähler potential as a remnant of the original gauge symmetry. To leading order, this gauge symmetry results in a shift symmetry for the Wilson line fields. As detailed in [21, 22], this interpretation is particularly natural in models which allow for a five-dimensional limit. To illustrate the role of the underlying Wilson line shift symmetry we follow these references and consider a four-dimensional  $\mathcal{N} = 1$  supersymmetric  $SU(5)$  GUT model that results from a five-dimensional  $SU(6)$  model compactified on an  $S^1$  orbifold.<sup>5</sup> Prior to orbifolding, the theory contains a four-dimensional chiral superfield  $\Phi$  in the adjoint of  $SU(6)$  whose real and imaginary part are, respectively, the scalar component of the five-dimensional vector multiplet and the Wilson line along  $S^1$ . The Wilson line shift symmetry, which descends from five-dimensional gauge invariance, ensures that the tree-level Kähler potential to quadratic order in  $\Phi$  takes the form

$$K = \frac{1}{2} \text{tr}(\Phi + \bar{\Phi})^2 f(S, \bar{S}) \quad (21)$$

with  $S$  collectively denoting other moduli fields. It is possible to choose the orbifold action in such a way that under the decomposition

$$SU(6) \rightarrow SU(5) \times U(1), \quad \mathbf{35} \rightarrow \mathbf{24} + \mathbf{1} + \mathbf{5} + \bar{\mathbf{5}} \quad (22)$$

only the component  $\mathbf{5} + \bar{\mathbf{5}}$  survives for  $\Phi$ . A suitable choice of discrete Wilson lines will lead to a breaking of  $SU(6) \rightarrow SU(4) \times SU(2) \times U(1)$  on one of the fixed points and thus

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<sup>5</sup>The assumption of  $SU(3) \times SU(2) \times U(1)$  gauge unification is only made for simplicity, and is not necessary in order to obtain Higgs fields in the correct representations from Wilson lines. On the contrary, in absence of low-energy supersymmetry the GUT idea is less compelling.

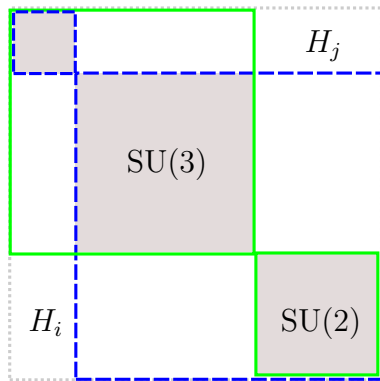


Figure 3: An illustration of the group theoretic origin of Higgs doublets from Wilson lines in the  $SU(6)$  orbifold case. The components of  $su(6)$  are displayed as a  $6 \times 6$  matrix. The generators of the gauge groups on the fixed points,  $SU(5) \times U(1)$  and  $SU(4) \times SU(2) \times U(1)$ , are marked by a dashed blue and solid green border respectively. The components corresponding to unbroken generators are shaded. The coset which corresponds to generators broken on both fixed points (and thus the massless components of  $\Phi$ ) is marked with  $H$ .

$SU(5) \rightarrow SU(3) \times SU(2) \times U(1)$ . This ensures that  $\Phi$  gives rise precisely to the chiral superfields  $H_u$  and  $H_d$  in the massless spectrum. For details of this projection see [31] and [21] for a 6d version and other variants. The situation is illustrated in Figure 3. In this case evaluation of (21) results in

$$K = \frac{1}{2} \text{tr}(\Phi + \bar{\Phi})^2 f(S, \bar{S}) = (H_u + \bar{H}_d)(\bar{H}_u + H_d) f(S, \bar{S}), \quad (23)$$

which is the starting point of our phenomenological analysis.

A similar structure of the Kähler potential should be encountered also in smooth heterotic compactifications and in Type II string constructions with D-branes.<sup>6</sup> However, a direct comparison between the orbifold point and smooth heterotic compactifications and furthermore with dual Type II models with branes is complicated. The intricate relation between heterotic orbifolds and smooth compactifications has been explored systematically only in the recent literature beginning with [28]. As pointed out above, the Kähler potential (20) can be corrected at subleading order (but still at string tree-level) by couplings involving pairs of twisted sector blow-up modes. Giving a non-zero vacuum expectation value to these fields smoothens out the orbifold into a heterotic compactification with vector bundles on an at least partially resolved Calabi-Yau space. In the presence of higher order terms involving these blow-up modes, the shift symmetry might well be broken and the structure (20) can be corrected.

With this in mind we take the established existence of the desired shift symmetry in heterotic orbifolds as an inspiration to search for similar shift symmetries in Type II models. We will make sure, though, to give arguments applicable entirely within Type II theory and independent of a possible duality to heterotic orbifolds.

<sup>6</sup>As is well known the target space dynamics (D-terms of anomalous  $U(1)$ s) drives heterotic orbifold models away from the orbifold point into the regime of compactification on at least partially resolved Calabi-Yau manifolds.

### 3.2 Shift symmetry for open string Wilson lines

At leading-order the Wilson line moduli  $\Phi^{(i)}$  associated with the theory on a (single) Type II D-brane enjoy a shift symmetry suitable for our purposes. For definiteness consider a single D6-brane in Type IIA string theory wrapping a special Lagrangian 3-cycle  $\Sigma$  with  $b_1(\Sigma)$  brane moduli. Each modulus is described by an  $\mathcal{N} = 1$  chiral superfield  $\Phi^{(i)}$  whose bosonic component  $\varphi^{(i)} + ia^{(i)}$  is the sum of a normal deformation  $\varphi^{(i)}$  and the Wilson line  $a^{(i)} = \int_{C_i} A$  (where  $C_i$  is one of the  $b_1(\Sigma)$  1-cycles on  $\Sigma$ ). The key observation is that in absence of charged matter states at the intersection of two D6-branes the Wilson line  $a^{(i)}$  does not couple non-derivatively in the effective action. This holds at tree-level in  $g_s$  and perturbatively in  $\alpha'$ . Let us recall the underlying CFT argument of [34], which generalises similar arguments from the heterotic string [33]: Non-derivative couplings of the Wilson line involve the zero-momentum limit ( $k = 0$ ) of the Wilson line vertex operator. In the 0-picture and in integrated form it is given by

$$V_{a^{(i)}}|_{k=0} = \int_{\partial D} A_\mu^{(i)}(X) \partial_\alpha X^\mu d\sigma^\alpha. \quad (24)$$

Here  $D$  is the open string worldsheet that appears in the scattering process, which at tree-level in  $g_s$  is topologically a disk, and  $A_\mu^{(i)}$  is polarised parallel to the brane. Since the Wilson line is flat,  $dA^{(i)} = 0$ , we can write  $A^{(i)} = d\chi^{(i)}$  if  $X(\partial D)$  is topologically trivial in  $\Sigma$ . Therefore the vertex operator vanishes by integration by parts [34]. Whenever the above reasoning goes through, no non-derivative terms involving the Wilson line are possible in the effective action. By holomorphicity this in particular excludes non-derivative terms involving the superfield  $\Phi^{(i)}$  in the superpotential. More importantly for us, it is also the origin of the shift-symmetric form of the brane-modulus Kähler potential.

The resulting shift symmetry is corrected at higher order: First, for topologically non-trivial  $X(\partial D)$ , as occur for scattering processes described by a worldsheet instanton, the above argument does not apply. This is in agreement with the non-perturbative breakdown of the classical shift-symmetry through superpotential terms for A-type brane moduli depending on  $e^{\Phi^{(i)}}$ . Second, if we compute a coupling involving, in addition to  $V_{a^{(i)}}$ , boundary changing vertex operators, the disk boundary is partitioned into several segments, and integration by parts need not give zero. Therefore non-derivative terms in the effective action involving the Wilson line fields and charged open string states located at the intersection of brane pairs are possible. This explains in particular the appearance of superpotential terms of the type  $W \supset \Phi^{(i)} \mathcal{O}$  with  $\mathcal{O}$  a product of charged localised states. Such terms likewise break the shift symmetry. Third, on higher genus worldsheets  $V_{a^{(i)}}|_{k=0}$  need not vanish either, resulting in string-loop corrections in  $g_s$  to the effective action which may involve contact terms in  $a^{(i)}$ .

Applying the above logic to the Kähler potential, we conclude that to second order in  $\Phi^{(i)}$  (and assuming just a single Wilson line for simplicity) it must exhibit a shift symmetric form of the type  $K = (\Phi + \bar{\Phi})^2 f(S, \bar{S})$ , where  $S$  now collectively denotes the closed string moduli of the compactification. In addition to this, subleading contributions to the Kähler potential exist which break the shift symmetry. As discussed they are

suppressed by  $g_s$ , non-perturbative in  $\Phi$  and  $\bar{\Phi}$  or possibly involve boundary changing charged fields. We will comment on the latter in the next section.

Note that these general assertions only depend on the structure of the open string CFT and must therefore hold independently of the chosen background.<sup>7</sup> Even though presented in the context of Type IIA D6-branes, the above worldsheet argument also governs the structure of Wilson line moduli on B-type branes, i.e. of D9, D7 or D5-branes in Type IIB orientifolds.

For D7 and D5-branes another set of moduli appears in the form of transverse deformation moduli. Being transversely polarised as opposed to parallel to the brane as in (24), no direct analogue of the CFT argument for Wilson lines can be made. This is in agreement with the generic appearance of a perturbative (in  $\alpha'$ ) superpotential for these deformation modes in presence of suitable fluxes. Similarly, it is a priori not clear that the Kähler potential exhibits a shift symmetry (at leading order in a suitable expansion) because the worldsheet instanton corrections of Type IIA generically map to perturbative corrections in Type IIB which spoil the shift symmetry.<sup>8</sup>

This identifies, at least in the supergravity limit of controllable worldsheet instanton corrections, the Wilson line moduli sector as a starting point to achieve a Kähler potential of the desired structure.

### 3.3 Shift symmetry for bulk matter

The analysis in the previous section has been for a single D-brane with gauge group  $U(1)$ . We now assess to what extent it is possible to generalise this to a leading-order shift symmetry in the Kähler potential of Wilson moduli transforming in the adjoint representation of the gauge group  $G$  along a stack of coincident D-branes. Upon gauge symmetry breaking such a symmetry induces a corresponding symmetry for those bulk matter states that descend from  $\Phi$ , similarly to the heterotic orbifold summarised in section 3.1. While a more quantitative analysis goes beyond the scope of this letter and is reserved for future work, we set out to describe the general picture.

Consider first a stack of Type IIA D6-branes carrying gauge group  $G$  and assume that the 3-cycle  $\Sigma$  within the Calabi-Yau 3-fold  $X_3$  wrapped by this stack admits one geometric modulus. The associated chiral superfield  $\Phi$  transforms in the adjoint of  $G$ . Under a breaking of the gauge group  $G \rightarrow H \times F$  the adjoint of  $G$  decomposes as

$$\text{ad}(G) \rightarrow (\text{ad}(H), 1) \oplus (1, \text{ad}(F)) \oplus \sum_i [(R_i, U_i) + c.c.], \quad (25)$$

where  $R_i$  and  $U_i$  are irreducible representations of  $H$  and  $F$ , respectively. Concretely we are interested in set-ups in which the surviving gauge group  $H$  either is directly

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<sup>7</sup>A systematic computation of the Kähler metric  $G_{\Phi\bar{\Phi}}$  in particular of untwisted matter fields on toroidal backgrounds has been performed in [35], in agreement with results found by duality with the heterotic string [36].

<sup>8</sup>This does not exclude the possible appearance of such shift symmetries for specific geometries or in particular regions of the moduli space [37].

$SU(2)_L \times U(1)_Y$  or contains it (in which case another breaking mechanism would have to be implemented in a second step). The gauge symmetry breaking must be such that the only surviving components of  $\Phi$  can be identified with  $H_u$  and  $H_d$ . The SM  $SU(3)$  factor can be associated with a different brane stack as is common in intersecting brane models, or it can be contained within  $H$  in a GUT construction.

The vertex operator argument presented around eq. (24) for a single D-brane makes use of the fact that the amplitude does not involve boundary changing operators. In particular, the Wilson line  $a^{(i)}$  itself is the excitation of an open string starting and ending on the same brane. For a stack of  $N$  coincident D-branes at generic position (i.e. not invariant under the orientifold projection) the original gauge group is  $G = U(N)$ . The Wilson line shift symmetry for  $N = 1$  as realized by a single brane directly carries over to a shift symmetry for the components of  $\Phi$  associated with each of the  $N$  Cartan generators  $U(1)^N \subset U(N)$ . These correspond to states from open strings starting and ending on one of the  $N$  branes within the stack. The shift symmetry of these fields together with the full  $U(N)$  gauge invariance constrains the possible couplings in a manner sufficient for our purposes:

Indeed, let  $\Phi = \Phi^a T^a$ , where  $T^a$  are the  $N^2$  generators of  $U(N)$ . We can restrict our attention to the  $N^2 - 1$  generators of  $SU(N) \subset U(N)$ . Within this subset of fields there is clearly no gauge invariant term linear in  $\Phi$ . At quadratic order, the most general Kähler potential can be built from the two independent sets of fields  $(\Phi + \bar{\Phi})^a$  and  $(\Phi - \bar{\Phi})^a$ , contracted in all possible ways with the unique invariant tensor,  $\delta^{ab}$ . Of the resulting three terms

$$(\Phi + \bar{\Phi})^a (\Phi + \bar{\Phi})^b \delta^{ab}, \quad (\Phi + \bar{\Phi})^a (\Phi - \bar{\Phi})^b \delta^{ab}, \quad (\Phi - \bar{\Phi})^a (\Phi - \bar{\Phi})^b \delta^{ab}, \quad (26)$$

the last two are forbidden for the Cartan generators and hence forbidden in general. This ensures the desired shift-symmetric structure of the quadratic-order Kähler potential.

At cubic order in  $\Phi$ , one can use the structure constants  $f^{abc}$  and build invariant expressions which can not be constrained using the vertex operator argument for the Cartan generators. Possible terms which violate the shift symmetry can be viewed as part of the corrections to the Kähler potential due to boundary changing operators. They are suppressed by an extra power of the Planck mass and do therefore not affect the phenomenological analysis of the previous section.

On top of that, the shift symmetry can be broken at tree-level by couplings to fields charged under the SM. The latter, however, are innocuous for our application because fields charged under the SM gauge group will have neither a bosonic VEV nor, generically, a non-trivial F-term. Thus such terms do not affect the structure of the Higgs mass matrix resulting from the tree-level Kähler potential.

The symmetry breaking (25) can be implemented by quotienting the compactification space  $X_3$  by a discrete symmetry group  $\mathbb{G}$  (which restricts to a symmetry of the 3-cycle  $\Sigma$  wrapped by the D6-brane) and suitably embedding the action of  $\mathbb{G}$  into  $G$ . In the case of a freely-acting symmetry group this is just the Hosotani-Witten mechanism of switching on discrete Wilson lines. Massless matter descending from  $\Phi$  is given by the zero modes of the Laplace operator acting on 1-forms on  $\Sigma$  twisted by the flat connection corresponding

to the Wilson line. This, however, leaves the components of  $\Phi$  in the adjoint of  $H$  massless because these correspond to the zero modes of the untwisted Laplace operator, which are unaffected. (Put differently, freely acting quotients do not reduce the fundamental group and so  $b_1(\Sigma)$  cannot decrease.)

What remains possible is to quotient  $X_3$  by a discrete symmetry group  $\mathbb{G}$  that restricts to a non-freely acting symmetry of  $\Sigma$ . This corresponds to an orbifold  $X/\mathbb{G}$  with non-trivial fixed-points, even though the covering space  $X_3$  is not necessarily toroidal. Nonetheless the string theory on  $X_3/\mathbb{G}$  is well-defined as long as  $\mathbb{G}$  is embedded into the holonomy group  $SU(3)$ . The cycle  $\Sigma$  wrapped by the D6-brane must pass through some of the fixed-points so that the brane is mapped to itself and the brane spectrum is projected. For a concise summary of this projection technique in the context of toroidal Type IIB branes see e.g. [36] and references therein, and the same methods apply to Type IIA branes.

Let us now comment on the situation for gauge theories on 7-branes in Type IIB orientifolds or F-theory. As reviewed at the end of section 3.2, there are two types of open string moduli of a 7-brane wrapping a holomorphic divisor  $S$ . The complex Wilson line moduli counted by elements in  $H^1(S)$  exhibit a shift symmetry, while for the geometric deformations given by<sup>9</sup>  $H^0(S, K_S) \simeq H^2(S, \mathcal{O})^*$  this symmetry will generically be broken in the manner discussed above.

If we start with a gauge group  $G$ , both types of fields give rise to chiral multiplets in the adjoint representation, which descend to matter in suitable representations upon breaking  $G \rightarrow H \times F$ . In addition to symmetry breaking via orbifolds, Type IIB models offer the possibility of considering non-flat gauge connections, i.e. non-trivial gauge flux. In the presence of gauge flux, the divisor  $S$  is endowed with a non-trivial holomorphic vector bundle  $L$ , whose structure group  $F$  is embedded into  $G$ . For simplicity we can focus on line bundles by taking  $F = U(1)$ . The massless  $\mathcal{N} = 1$  chiral superfields in representation  $R_i$  with  $U(1)$  charge  $q_i$  are given by the zero modes of the twisted Laplace operator acting on one- and two-forms. Concretely, the chiral bulk spectrum is given by

$$H^1(S, L^{q_i}) \oplus H^2(S, L^{-q_i}). \quad (27)$$

The chiral superfields in  $\bar{R}_i$  are counted by

$$H^1(S, L^{-q_i}) \oplus H^2(S, L^{q_i}). \quad (28)$$

See e.g. [32, 38] for a recent discussion of details of the matter spectrum on 7-branes. It is tempting to ascribe to the elements in  $H^1(S, L^{q_i})$  a Wilson line shift symmetry. However, unlike in the case of an orbifold, these states are *not* directly related to a “universal” Wilson line in the adjoint of the underlying gauge group  $G$ , which would be counted by  $H^1(S, \mathcal{O})$ . It is therefore not obvious that the shift symmetry argument for such an underlying Wilson line implies a Kähler potential of the form (23).<sup>10</sup>

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<sup>9</sup>Here  $K_S$  is the canonical bundle of  $S$  and  $\mathcal{O}$  the trivial bundle.

<sup>10</sup> Indeed, the CFT argument for the shift symmetry only holds for flat Wilson lines. Note, however, that the vector bundles on a B-type brane are holomorphic, i.e. their curvature satisfies  $F^{2,0} = 0 = F^{0,2}$ . While this implies that curvature is non-zero in certain directions, one may hope that others are still protected from non-derivative couplings.

Some helpful light can be shed on this question by comparison with heterotic orbifolds. The blow-up of such orbifolds corresponds to a smooth heterotic compactification with non-trivial vector bundles [28]. Both the twisted and the untwisted matter maps to different elements of the same cohomology group on the resolved space (see also the discussion in the appendix of [29]). This suggests that one combination of these fields *can* be viewed as the analogue of orbifold bulk matter. We reiterate, however, that the resolution process might significantly correct the shift symmetric Kähler potential due to higher-order terms involving the blow-up modes. We leave it for future work to study the shift symmetries of bulk matter in presence of vector bundles.

Models with bulk matter fields have recently been re-addressed also in the context of F-theory GUT model building. Compatibility of a bulk Higgs field with the desired Yukawa structure often requires non-perturbative effects in form of exceptional singularity enhancements as considered recently in [39]. The construction of concrete setups which in particular accomodate the SM Yukawa couplings is an exciting challenge which we plan to address in the near future.

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